WHY DO WE NEED NON-BARYONIC DARK MATTER? [1]

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Abstract

Observational evidence along with theoretical arguments which call for non-baryonic dark matter are reviewed. A brief summary of the dark matter session is included.

There is increasing evidence that relative to the visible matter in the Universe, which is in the form of baryons, there is considerably more matter in the Universe that we don't see [1]. Here, I will review some of the motivations for dark matter in the Universe. The best observational evidence is found on the scale of galactic halos and comes from the observed flat rotation curves of galaxies. There is also mounting evidence for dark matter in elliptical galaxies as well as clusters of galaxies coming from X-ray observations of these objects. Also, direct evidence has been obtained through the study of gravitational lenses. In theory, we believe there is much more matter because 1) inflation tells us so (and there is at present no good alternative to inflation) and 2) our current understanding of galaxy formation only makes sense if there is more matter than we see. One can also make a strong case for the existence of non-baryonic dark matter in particular. The recurrent problem with baryonic dark matter is that not only is it very difficult to hide baryons, but the standard model of primordial nucleosynthesis would have to discarded if all of the dark matter is baryonic. Fortunately, as will be covered at length in these proceedings, there are several attractive alternatives to baryonic dark matter.

Before embarking on the subject of dark matter, it will be useful to review the relevant quantities from the standard big bang model. In a Friedmann-Robertson-Walker Universe, the expansion rate of the Universe (the Hubble parameter) is related to the energy density ρ and curvature constant k by

$$H^2 = \left(\frac{\dot{R}}{R}\right)^2 = \frac{8\pi G}{3}\rho - \frac{k}{R^2} \tag{1}$$

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assuming no cosmological constant, where $k = \pm 1, 0$ for a closed, open or spatially flat Universe, and R is the cosmological scale factor. When $k = 0$, the energy density takes its "critical" value, $\rho = \rho_c = 3H^2/8\pi G = 1.88 \times 10^{-29} h_o^2$ g cm^{-3} where $h_o = H_o/100$ km s^{-1} Mpc^{-1} is the scaled present value of the Hubble parameter. The cosmological density parameter is defined by $\Omega = \rho/\rho_c$ and by rewriting eq. (1) we can relate k to Ω and H by

$$\frac{k}{R^2} = (\Omega - 1)H^2 \tag{2}$$

so that $k = +1, -1, 0$ corresponds to $\Omega > 1, < 1, = 1$.

In very broad terms, observational limits on the cosmological parameters are: $0.2 \lesssim \Omega \lesssim 2$ and $0.4 \lesssim h_o \lesssim 1.0$ [2]. The cosmological density is however sensitive to the particular scale being observed (at least on small scales). Accountably visible matter contributes in total only a small fraction to the overall density, giving $\Omega_V \sim .003 - .01$. In the bright central parts of galaxies, the density is larger $\Omega \sim 0.02 - 0.1$. On larger scales, that of binaries and small groups of galaxies, $\Omega \simeq 0.05 - 0.3$. On even larger scales the density may be large enough to support $\Omega \simeq 1.0$. Though there are no astronomical observations to support $\Omega > 1$, limits based on the deceleration of the Universe only indicate [2] that $\Omega \lesssim 2$.

The age of the Universe is also very sensitive to these parameters. Again, in the absence of a cosmological constant we have,

$$H_o t_U = \int_0^1 (1 - \Omega + \Omega/x)^{-1/2} dx \tag{3}$$

For $t_U > 13$ Gyr, $\Omega h_o^2 < 0.25$ if $h_o > 0.5$ and $\Omega h_o^2 < 0.45$ if $h_o > 0.4$. While for $t_U > 10$ Gyr, $\Omega h_o^2 < 0.8$ if $h_o > 0.5$ and $\Omega h_o^2 < 1.1$ if $h_o > 0.4$.

There is, in fact, good evidence for dark matter on the scale of galaxies (and their halos). Assuming that galaxies are in virial equilibrium, one expects that by Newton's Laws one can relate the mass at a given distance r, from the center of a galaxy to its rotational velocity

$$M(r) \propto v^2 r/G_N \tag{4}$$

The rotational velocity, v, is measured [3, 4, 5] by observing 21 cm emission lines in HI regions (neutral hydrogen) beyond the point where most of the light in the galaxy ceases. A compilation of nearly 1000 rotation curves of spiral galaxies have been plotted in [6] as a function of r for varying brightnesses. If the bulk of the mass is associated with light, then beyond the point where most of the light stops M would be constant and $v^2 \propto 1/r$. This is not the case, as the rotation curves appear to be flat, i.e., $v \sim$ constant outside the core of the galaxy. This implies that $M \propto r$ beyond the point where the light stops. This is one of the strongest pieces of evidence for the existence of dark matter. Velocity measurements indicate dark matter in elliptical galaxies as well [7].

Galactic rotation curves are not the only observational indication for the existence of dark matter. X-ray emitting hot gas in elliptical galaxies also provides an important piece of evidence for dark matter. As an example, consider the large elliptical M87. The detailed profiles of the temperature and density of the hot X-ray emitting gas have been mapped out [8]. By assuming hydrostatic equilibrium, these measurements allow one to determine the overall mass distribution in the galaxy necessary to bind the hot gas. Based on an isothermal model with temperature $kT = 3$keV (which leads to a conservative estimate of the total mass),

Fabricant and Gorenstein [8] predicted that the total mass out to a radial distance of 392 Kpc, is $5.7 \times 10^{13} M_\odot$ whereas the mass in the hot gas is only $2.8 \times 10^{12} M_\odot$ or only 5% of the total. The visible mass is expected to contribute only 1% of the total. The inferred value of Ω based on M87 would be ~ 0.2.

M87 is not the only example of an elliptical galaxy in which X-ray emitting hot gas is observed to indicate the presence of dark matter. At this meeting, Foreman [9], showed several examples of ellipiticals with large mass to light ratios. For example in the case of N4472, while the optical observations go out to 25 kpc, the X-ray gas is seen out to 75 kpc, indicating M/L's of about 60 at 70 kpc and up to 90 at 100 kpc. Similar inferences regarding the existence of dark matter can be made from the X-ray emission from small groups of galaxies [10, 11].

On very large scales, it is possible to get an estimate of Ω from the distribution of peculiar velocities. On scales, λ, where perturbations, δ, are still small, peculiar velocities can be expressed [12] as $v \sim H\lambda\delta\Omega^{0.6}$. On these scales, measurements of the peculiar velocity field from the IRAS galaxy catalogue indicate that indeed Ω is close to unity [13]. Another piece of evidence on large scales, is available from gravitational lensing [14]. The systematic lensing of the roughly 150,000 galaxies per $\deg^2$ at redshifts between $z = 1 - 3$ into arcs and arclets allow one to trace the matter distribution in a foreground cluster. Van Waerbeke discussed recent results of weak gravitational lensing looking at systems, which if virialized, have mass to light ratios in the range 400–1000 and correspond to values of Ω between 0.25 and 0.6 [15]. Unfortunately, none of these observations reveal the identity of the dark matter.

Theoretically, there is no lack of support for the dark matter hypothesis. The standard big bang model including inflation almost requires that $\Omega = 1$ [16]. The simple and unfortunate fact that at present we do not even know whether Ω is larger or smaller than one indicates that we do not know the sign of the curvature term, further implying that it is subdominant in Eq. (1)

$$\frac{k}{R^2} < \frac{8\pi G}{3}\rho \tag{5}$$

In an adiabatically expanding Universe, $R \sim T^{-1}$ where T is the temperature of the thermal photon background. Therefore the quantity

$$\hat{k} = \frac{k}{R^2 T^2} < \frac{8\pi G}{3 T_o^2} < 2 \times 10^{-58} \tag{6}$$

is dimensionless and constant in the standard model. The smallness of $\hat{k}$ is known as the curvature problem and can be resolved by a period of inflation. Before inflation, let us write $R = R_i$, $T = T_i$ and $R \sim T^{-1}$. During inflation, $R \sim T^{-1} \sim e^{Ht}$, where H is constant. After inflation, $R = R_f \gg R_i$ but $T = T_f = T_R \lesssim T_i$ where T_R is the temperature to which the Universe reheats. Thus $R \not\sim T^{-1}$ and $\hat{k} \to 0$ is not constant. But from Eqs. (2) and (6) if $\hat{k} \to 0$ then $\Omega \to 1$, and since typical inflationary models contain much more expansion than is necessary, Ω becomes exponentially close to one.

If this is the case and $\Omega = 1$, then we know two things: Dark matter exists, since we don't see $\Omega = 1$ in luminous objects, and most (at least 90%) of the dark matter is not baryonic. The latter conclusion is a result from big bang nucleosynthesis [17, 18], which constrains the baryon-to-photon ratio $\eta = n_B/n_\gamma$ to

$$1.4 \times 10^{-10} < \eta < 3.8 \times 10^{-10} \tag{7}$$

which corresponds to a limit on Ω_B

$$0.005 < \Omega_B < 0.09 \tag{8}$$

for $0.4 \lesssim h_o \lesssim 1.0$. Thus $1 - \Omega_B$ is not only dark but also non-baryonic. I will return to big bang nucleosynthesis below.

Another important piece of theoretical evidence for dark matter comes from the simple fact that we are here living in a galaxy. The type of perturbations produced by inflation [19] are, in most models, adiabatic perturbations ($\delta\rho/\rho \propto \delta T/T$), and I will restrict my attention to these. Indeed, the perturbations produced by inflation also have the very nearly scale-free spectrum described by Harrison and Zeldovich [20]. When produced, scale-free perturbations fall off as $\frac{\delta\rho}{\rho} \propto l^{-2}$ (increase as the square of the wave number). At early times $\delta\rho/\rho$ grows as t until the time when the horizon scale (which is proportional to the age of the Universe) is comparable to l. At later times, the growth halts (the mass contained within the volume l^3 has become smaller than the Jean's mass) and $\frac{\delta\rho}{\rho} = \delta$ (roughly) independent of the scale l. When the Universe becomes matter dominated, the Jean's mass drops dramatically and growth continues as $\frac{\delta\rho}{\rho} \propto R \sim 1/T$. For an overview of the evolution of density perturbations and the resulting observable spectrum see [21]. The transition to matter dominance is determined by setting the energy densities in radiation (photons and any massless neutrinos) equal to the energy density in matter (baryons and any dark matter). For three massless neutrinos and baryons (no dark matter), matter dominance begins at

$$T_m = 0.22 m_B \eta \tag{9}$$

and for $\eta < 3.8 \times 10^{-10}$, this corresponds to $T_m < 0.08$ eV.

The subsequent non-linear growth in $\delta\rho/\rho$ was discussed in these sessions in some detail at this meeting. In particular, there was a considerable discussion of the effects of the non-linear regime on the power spectrum and the appearance of non-Gaussian features such as skewness and kurtosis [22]. Colombi [23] discussed the 2- and 3-point correlation functions. Numerous simulations were presented to reflect the hydrodynamics of galaxy formation [24] and the origin of the large scale bias [25].

Because we are considering adiabatic perturbations, there will be anisotropies produced in the microwave background radiation on the order of $\delta T/T \sim \delta$. The value of δ, the amplitude of the density fluctuations at horizon crossing, has now been determined by COBE [26], $\delta = (5.7 \pm 0.4) \times 10^{-6}$. Without the existence of dark matter, $\delta\rho/\rho$ in baryons could then achieve a maximum value of only $\delta\rho/\rho \sim A_\lambda \delta(T_m/T_o) \lesssim 2 \times 10^{-3} A_\lambda$, where $T_o = 2.35 \times 10^{-4}$ eV is the present temperature of the microwave background and $A_\lambda \sim 1 - 10$ is a scale dependent growth factor. The overall growth in $\delta\rho/\rho$ is too small to argue that growth has entered a nonlinear regime needed to explain the large value (10^5) of $\delta\rho/\rho$ in galaxies.

Dark matter easily remedies this dilemma in the following way. The transition to matter dominance is determined by setting equal to each other the energy densities in radiation (photons and any massless neutrinos) and matter (baryons and any dark matter). While if we suppose that there exists dark matter with an abundance $Y_\chi = n_\chi/n_\gamma$ (the ratio of the number density of χ's to photons) then

$$T_m = 0.22 m_\chi Y_\chi \tag{10}$$

Since we can write $m_\chi Y_\chi/\text{GeV} = \Omega_\chi h^2/(4 \times 10^7)$, we have $T_m/T_o = 2.4 \times 10^4 \Omega_\chi h^2$ which is to be compared with 350 in the case of baryons alone. The baryons, although still bound

to the radiation until decoupling, now see deep potential wells formed by the dark matter perturbations to fall into and are no longer required to grow at the rate $\delta\rho/\rho \propto R$.

With regard to dark matter and galaxy formation, all forms of dark matter are not equal. They can be distinguished by their relative temperature at T_m [27]. Particles which are still largely relativistic at T_m (like neutrinos or other particles with $m_\chi < 100$ eV) have the property [28] that (due to free streaming) they erase perturbations out to very large scales given by the Jean's mass

$$M_J = 3 \times 10^{18} \frac{M_\odot}{m_\nu{}^2(eV)} \tag{11}$$

Thus, very large scale structures form first and galaxies are expected to fragment out later. Particles with this property are termed hot dark matter particles. Cold particles ($m_\chi > 1$ MeV) have the opposite behavior. Small scale structure forms first aggregating to form larger structures later. Neither candidate is completely satisfactory when the resulting structure is compared to the observations. For more details, I refer the reader to reviews in refs. [1].

The most promising possibility we have for unscrambling the various possible scenarios for structure formation is the careful analysis of the observed power spectrum. Rewriting the density contrast in Fourier space,

$$\delta(k) \propto \int d^3x \frac{\delta\rho}{\rho}(x)e^{ik\cdot x} \tag{12}$$

the power spectrum is just

$$P(k) = \langle|\delta(k)|^2\rangle \tag{13}$$

which is often written in term of a transfer function and a power law, $P \sim T(k)k^n$. (The flat spectrum produced by inflation has $n = 1$.) The data contributing to $P(k)$ include observations of galaxy distributions, velocity distributions and of course the CMB [29]. However to make a comparison with our theoretical expectations, we rely on n-body simulations and a deconvolution of the theory from the observations. Overall, there is actually a considerable amount of concordance with our expectations. The velocity distributions indicate that $0.3 < \Omega < 1$ and the power spectrum is roughly consistent with an $\Omega = 1$, and $n = 1$, cold dark matter Universe.

Future surveys [30, 31] will however, most certainly dramatically improve our understanding of the detailed features of the power spectrum and their theoretical interpretations. We should in principle be able to distinguish between a mixed dark matter and a cold dark matter $\Omega = 1$ Universe, whether or not $\Omega < 1$, with CDM, the presence of a cosmological constant, or a tilt in the spectrum. Zurek [32], stressed the importance of numerical simulations in this context. Future probes of the small scale anisotropy [31, 33] should in addition be able to determine the values of not only Ω, but also Ω_B and h_o to a high degree of accuracy through a careful analysis of the Doppler peak in the power spectrum.

Accepting the dark matter hypothesis, the first choice for a candidate should be something we know to exist, baryons. Though baryonic dark matter can not be the whole story if $\Omega = 1$, the identity of the dark matter in galactic halos, which appear to contribute at the level of $\Omega \sim 0.05$, remains an important question needing to be resolved. A baryon density of this magnitude is not excluded by nucleosynthesis. Indeed we know some of the baryons are dark since $\Omega \lesssim 0.01$ in the disk of the galaxy.

It is quite difficult, however, to hide large amounts of baryonic matter [34]. Sites for halo baryons that have been discussed include snowballs, which tend to sublimate, cold hydrogen

gas, which must be supported against collapse, and hot gas, which can be excluded by the X-ray background. Stellar objects (collectively called MACHOs for macroscopic compact halo objects) must either be so small ($M < 0.08\ M_\odot$) so as not to have begun nuclear burning or so massive so as to have undergone total gravitational collapse without the ejection of heavy elements. Intermediate mass stars are generally quite problematic because either they are expected to still reside on the main-sequence today and hence would be visible, or they would have produced an excess of heavy elements.

On the other hand, MACHOs are a candidate which are testable by the gravitational microlensing of stars in a neighboring galaxy such as the LMC [35]. By observing millions of stars and examining their intensity as a function of time, it may be possible to determine the presence of dark objects in our halo. It is expected that during a lensing event, a star's intensity will rise in an achromatic fashion over a period $\delta t \sim 3\ \sqrt{M/.001 M_\odot}$ days. Indeed, microlensing candidates have been found [36]. For low mass objects, those with $M < 0.1 M_\odot$, it appears however that the halo fraction of Machos is small, $\approx 0.19^{+.16}_{-.10}$ [37]. There have been some recent reports of events with longer duration [38] leading to the speculation of a white dwarf population in the halo. Though it is too early to determine the implications of these observations, they are very encouraging in that perhaps this issue can and will be decided.

The degree to which baryons can contribute to dark matter depends ultimately on the overall baryon contribution to Ω which is constrained by nucleosynthesis. Because of its importance regarding the issue of dark matter and in particular non-baryonic dark matter, I want to review the status of big bang nucleosynthesis.

Conditions for the synthesis of the light elements were attained in the early Universe at temperatures $T \lesssim 1$ MeV, corresponding to an age of about 1 second. Given a single input parameter, the baryon-to-photon ratio, η, the theory is capable of predicting the primordial abundances of D/H, ^{3}He/H, ^{4}He/H and ^{7}Li/H. The comparison of these predictions with the observational determination of the abundances of the light elements not only tests the theory but also fixes the value of η.

The dominant product of big bang nucleosynthesis is ^{4}He resulting in an abundance of close to 25 % by mass. In the standard model, the ^{4}He mass fraction depends only weakly on η. When we go beyond the standard model, the ^{4}He abundance is very sensitive to changes in the expansion rate which can be related to the effective number of neutrino flavors. Lesser amounts of the other light elements are produced: D and ^{3}He at the level of about 10^{-5} by number, and ^{7}Li at the level of 10^{-10} by number.

There is now a good collection of abundance information on the ^{4}He mass fraction, Y, O/H, and N/H in over 50 extragalactic HII (ionized hydrogen) regions [39, 40, 41]. The observation of the heavy elements is important as the helium mass fraction observed in these HII regions has been augmented by some stellar processing, the degree to which is given by the oxygen and nitrogen abundances. In an extensive study based on the data in [39, 40], it was found [42] that the data is well represented by a linear correlation for Y vs. O/H and Y vs. N/H. It is then expected that the primordial abundance of ^{4}He can be determined from the intercept of that relation. The overall result of that analysis indicated a primordial mass fraction, $Y_p = 0.232 \pm 0.003$. In [43], the stability of this fit was verified by a Monte-Carlo analysis showing that the fits were not overly sensitive to any particular HII region. In addition, the data from [41] was also included, yielding a ^{4}He mass fraction [43]

$$Y_p = 0.234 \pm 0.003 \pm 0.005 \tag{14}$$

The second uncertainty is an estimate of the systematic uncertainty in the abundance determination.

The ^{7}Li abundance is also reasonably well known. In old, hot, population-II stars, ^{7}Li is found to have a very nearly uniform abundance. For stars with a surface temperature $T > 5500$ K and a metallicity less than about 1/20th solar (so that effects such as stellar convection may not be important), the abundances show little or no dispersion beyond that which is consistent with the errors of individual measurements. Indeed, as detailed in [44], much of the work concerning ^{7}Li has to do with the presence or absence of dispersion and whether or not there is in fact some tiny slope to a [Li] = log ^{7}Li/H + 12 vs. T or [Li] vs. [Fe/H] relationship. There is ^{7}Li data from nearly 100 halo stars, from a variety of sources. I will use the value given in [45] as the best estimate for the mean ^{7}Li abundance and its statistical uncertainty in halo stars

$$\text{Li/H} = (1.6 \pm 0.1^{+0.4+1.6}_{-0.3-0.5}) \times 10^{-10} \tag{15}$$

The first error is statistical, and the second is a systematic uncertainty that covers the range of abundances derived by various methods. The third set of errors in Eq. (15) accounts for the possibility that as much as half of the primordial ^{7}Li has been destroyed in stars, and that as much as 30% of the observed ^{7}Li may have been produced in cosmic ray collisions rather than in the Big Bang. Observations of ^{6}Li, Be, and B help constrain the degree to which these effects play a role [46]. For ^{7}Li, the uncertainties are clearly dominated by systematic effects.

Turning to D/H, we have three basic types of abundance information: 1) ISM data, 2) solar system information, and perhaps 3) a primordial abundance from quasar absorption systems. The best measurement for ISM D/H is [47]

$$(\text{D/H})_{\text{ISM}} = 1.60 \pm 0.09^{+0.05}_{-0.10} \times 10^{-5} \tag{16}$$

This value may not be universal (or galactic as the case may be) since there may be some real dispersion of D/H in the ISM [48]. The solar abundance of D/H is inferred from two distinct measurements of ^{3}He. The solar wind measurements of ^{3}He as well as the low temperature components of step-wise heating measurements of ^{3}He in meteorites yield the presolar (D + ^{3}He)/H ratio, as D was efficiently burned to ^{3}He in the Sun's pre-main-sequence phase. These measurements indicate that [49, 50]

$$\left(\frac{\text{D} + {}^3\text{He}}{\text{H}}\right)_\odot = (4.1 \pm 0.6 \pm 1.4) \times 10^{-5} \tag{17}$$

The high temperature components in meteorites are believed to yield the true solar ^{3}He/H ratio of [49, 50]

$$\left(\frac{{}^3\text{He}}{\text{H}}\right)_\odot = (1.5 \pm 0.2 \pm 0.3) \times 10^{-5} \tag{18}$$

The difference between these two abundances reveals the presolar D/H ratio, giving,

$$(\text{D/H})_\odot \approx (2.6 \pm 0.6 \pm 1.4) \times 10^{-5} \tag{19}$$

Finally, there have been several recent reported measurements of D/H is high redshift quasar absorption systems. Such measurements are in principle capable of determining the primordial value for D/H and hence η, because of the strong and monotonic dependence of D/H on η. However, at present, detections of D/H using quasar absorption systems indicate both a high

and low value of D/H. As such, we caution that these values may not turn out to represent the true primordial value. The first of these measurements [51] indicated a rather high D/H ratio, $D/H \approx 1.9 - 2.5 \times 10^{-4}$. A recent re-observation of the high D/H absorption system has been resolved into two components, both yielding high values with an average value of $D/H = 1.9^{+0.4}_{-0.3} \times 10^{-4}$ [52] as well as an additional system with a similar high value [53]. Other high D/H ratios were reported in [54]. However, there are reported low values of D/H in other such systems [55] with values $D/H \simeq 2.5 \times 10^{-5}$, significantly lower than the ones quoted above. It is probably premature to use this value as the primordial D/H abundance in an analysis of big bang nucleosynthesis, but it is certainly encouraging that future observations may soon yield a firm value for D/H. It is however important to note that there does seem to be a trend that over the history of the Galaxy, the D/H ratio is decreasing, something we expect from galactic chemical evolution. Of course the total amount of deuterium astration that has occurred is still uncertain, and model dependent.

There are also several types of ^{3}He measurements. As noted above, meteoritic extractions yield a presolar value for ^{3}He/H as given in Eq. (18). In addition, there are several ISM measurements of ^{3}He in galactic HII regions [56] which also show a wide dispersion

$$\left(\frac{^3\mathrm{He}}{\mathrm{H}}\right)_{\mathrm{HII}} \simeq 1 - 5 \times 10^{-5} \tag{20}$$

There is also a recent ISM measurement of ^{3}He [57] with

$$\left(\frac{^3\mathrm{He}}{\mathrm{H}}\right)_{\mathrm{ISM}} = 2.1^{+.9}_{-.8} \times 10^{-5} \tag{21}$$

Finally there are observations of ^{3}He in planetary nebulae [58] which show a very high ^{3}He abundance of ^{3}He/H $\sim 10^{-3}$.

Each of the light element isotopes can be made consistent with theory for a specific range in η. Overall consistency of course requires that the range in η agree among all four light elements. ^{3}He (together with D) has stood out in its importance for BBN, because it provided a (relatively large) lower limit for the baryon-to-photon ratio [59], $\eta_{10} > 2.8$. This limit for a long time was seen to be essential because it provided the only means for bounding η from below and in effect allows one to set an upper limit on the number of neutrino flavors [60], N_ν, as well as other constraints on particle physics properties. That is, the upper bound to N_ν is strongly dependent on the lower bound to η. This is easy to see: for lower η, the ^{4}He abundance drops, allowing for a larger N_ν, which would raise the ^{4}He abundance. However, for $\eta < 4 \times 10^{-11}$, corresponding to $\Omega h^2 \sim .001 - .002$, which is not too different from galactic mass densities, there is no bound whatsoever on N_ν [61]. Of course, with the improved data on ^{7}Li, we do have lower bounds on η which exceed 10^{-10}.

In [59], it was argued that since stars (even massive stars) do not destroy ^{3}He in its entirety, we can obtain a bound on η from an upper bound to the solar D and ^{3}He abundances. One can in fact limit [59, 62] the sum of primordial D and ^{3}He by applying the expression below at $t = \odot$

$$\left(\frac{\mathrm{D} + {}^3\mathrm{He}}{\mathrm{H}}\right)_{\mathrm{p}} \leq \left(\frac{\mathrm{D}}{\mathrm{H}}\right)_{\mathrm{t}} + \frac{1}{g_3}\left(\frac{^3\mathrm{He}}{\mathrm{H}}\right)_{\mathrm{t}} \tag{22}$$

In (22), g_3 is the fraction of a star's initial D and ^{3}He which survives as ^{3}He. For $g_3 > 0.25$ as suggested by stellar models, and using the solar data on D/H and ^{3}He/H, one finds $\eta_{10} > 2.8$.

The limit $\eta_{10} > 2.8$ derived using (22) is really a one shot approximation. Namely, it is assumed that material passes through a star no more than once. To determine whether or not the solar (and present) values of D/H and ^{3}He/H can be matched it is necessary to consider models of galactic chemical evolution [63]. In the absence of stellar ^{3}He production, particularly by low mass stars, it was shown [64] that there are indeed suitable choices for a star formation rate and an initial mass function to: 1) match the D/H evolution from a primordial value $(D/H)_p = 7.5 \times 10^{-5}$, corresponding to $\eta_{10} = 3$, through the solar and ISM abundances, while 2) at the same time keeping the ^{3}He/H evolution relatively flat so as not to overproduce ^{3}He at the solar and present epochs. This was achieved for $g_3 > 0.3$. Even for $g_3 \sim 0.7$, the present ^{3}He/H could be matched, though the solar value was found to be a factor of 2 too high. For $(D/H)_p \simeq 2 \times 10^{-4}$, corresponding to $\eta_{10} \simeq 1.7$, models could be found which destroy D sufficiently; however, overproduction of ^{3}He occurred unless g_3 was tuned down to about 0.1.

In the context of models of galactic chemical evolution, there is, however, little justification a priori for neglecting the production of ^{3}He in low mass stars. Indeed, stellar models predict that considerable amounts of ^{3}He are produced in stars between 1 and 3 $M_\odot$. For $M < 8M_\odot$, Iben and Truran [65] calculate

$$(^3\text{He/H})_f = 1.8 \times 10^{-4} \left(\frac{M_\odot}{M}\right)^2 + 0.7 \left[(D + \ ^3\text{He})/\text{H}\right]_i \qquad (23)$$

so that at $\eta_{10} = 3$, and $((D + \ ^3\text{He})/\text{H})_i = 9 \times 10^{-5}$, $g_3(1M_\odot) = 2.7$! It should be emphasized that this prediction is in fact consistent with the observation of high ^{3}He/H in planetary nebulae [58].

Generally, implementation of the ^{3}He yield in Eq. (23) in chemical evolution models leads to an overproduction of ^{3}He/H particularly at the solar epoch [66, 67]. In Figure 1, the evolution of D/H and ^{3}He/H is shown as a function of time from [49, 66]. The solid curves show the evolution in a simple model of galactic chemical evolution with a star formation rate proportional to the gas density and a power law IMF (see [66]) for details). The model was chosen to fit the observed deuterium abundances. However, as one can plainly see, ^{3}He is grossly overproduced (the deuterium data is represented by squares and ^{3}He by circles). Depending on the particular model chosen, it may be possible to come close to at least the upper end of the range of the ^{3}He/H observed in galactic HII regions [56], although the solar value is missed by many standard deviations.

The overproduction of ^{3}He relative to the solar meteoritic value seems to be a generic feature of chemical evolution models when ^{3}He production in low mass stars is included. In [49], a more extreme model of galactic chemical evolution was tested. There, it was assumed that the initial mass function was time dependent in such a way so as to favor massive stars early on (during the first two Gyr of the galaxy). Massive stars are preferential from the point of view of destroying ^{3}He. However, massive stars are also proficient producers of heavy elements and in order to keep the metallicity of the disk down to acceptable levels, supernovae driven outflow was also included. The degree of outflow was limited roughly by the observed metallicity in the intergalactic gas in clusters of galaxies. One further assumption was necessary; we allowed the massive stars to lose their ^{3}He depleted hydrogen envelopes prior to explosion. Thus only the heavier elements were expulsed from the galaxy. With all of these (semi-defensible) assumptions, ^{3}He was still overproduced at the solar epoch, as shown by the dashed curve in Figure 1. Though there certainly is an improvement in the evolution of ^{3}He without reducing the yields of low mass stars, it is hard to envision much further reduction in the solar ^{3}He
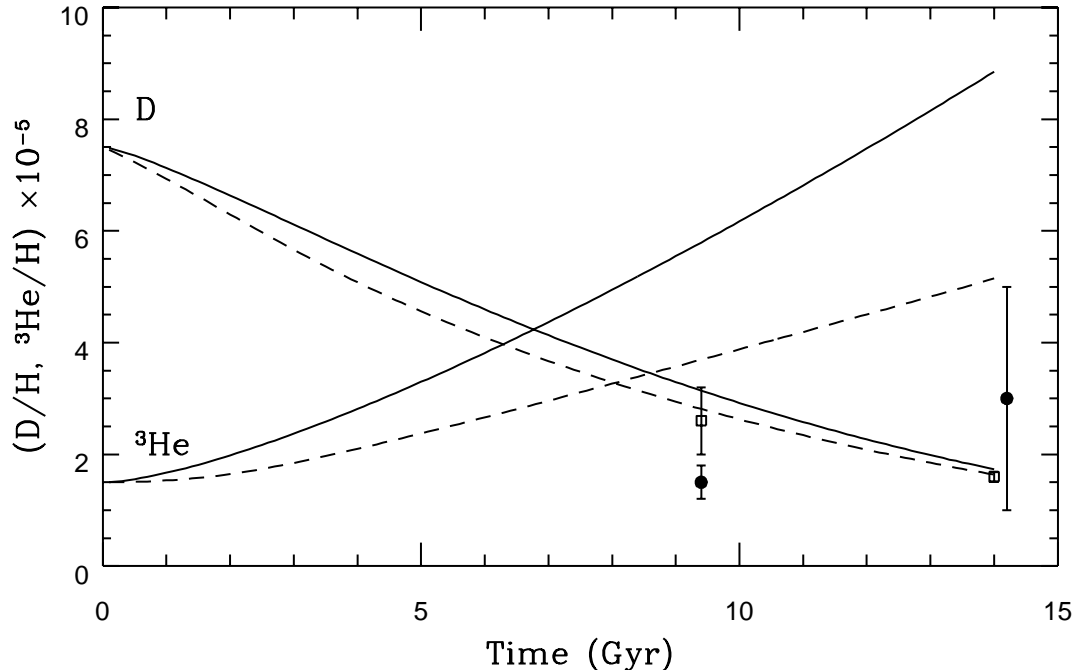
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Figure 1: The evolution of D and ^{3}He with time.

predicted by these models. The only conclusion that we can make at this point is that there is clearly something wrong with our understanding of ^{3}He in terms of either chemical evolution, stellar evolution or perhaps even the observational data.

Given the magnitude of the problems concerning ^{3}He, it would seem unwise to make any strong conclusion regarding big bang nucleosynthesis which is based on ^{3}He. Perhaps as well some caution is deserved with regard to the recent D/H measurements, although if the present trend continues and is verified in several different quasar absorption systems, then D/H will certainly become our best measure for the baryon-to-photon ratio. Given the current situation however, it makes sense to take a step back and perform an analysis of big bang nucleosynthesis in terms of the element isotopes that are best understood, namely ^{4}He and ^{7}Li.

Monte Carlo techniques are proving to be a useful form of analysis for big bang nucleosynthesis [68, 69]. In [18], we performed just such an analysis using only ^{4}He and ^{7}Li. It should be noted that in principle, two elements should be sufficient for constraining the one parameter (η) theory of BBN. We begin by establishing likelihood functions for the theory and observations. For example, for ^{4}He, the theoretical likelihood function takes the form

$$L_{\mathrm{BBN}}(Y, Y_{\mathrm{BBN}}) = e^{-(Y-Y_{\mathrm{BBN}}(\eta))^2/2\sigma_1^2} \qquad (24)$$

where $Y_{\mathrm{BBN}}(\eta)$ is the central value for the ^{4}He mass fraction produced in the big bang as predicted by the theory at a given value of η, and σ_1 is the uncertainty in that value derived from the Monte Carlo calculations [69] and is a measure of the theoretical uncertainty in the big bang calculation. Similarly one can write down an expression for the observational likelihood function. In this case we have two sources of errors, as discussed above, a statistical uncertainty, σ_2 and a systematic uncertainty, σ_{sys}. Here, I will assume that the systematic error is described by a top hat distribution [69, 70]. The convolution of the top hat distribution and the Gaussian (to describe the statistical errors in the observations) results in the difference of
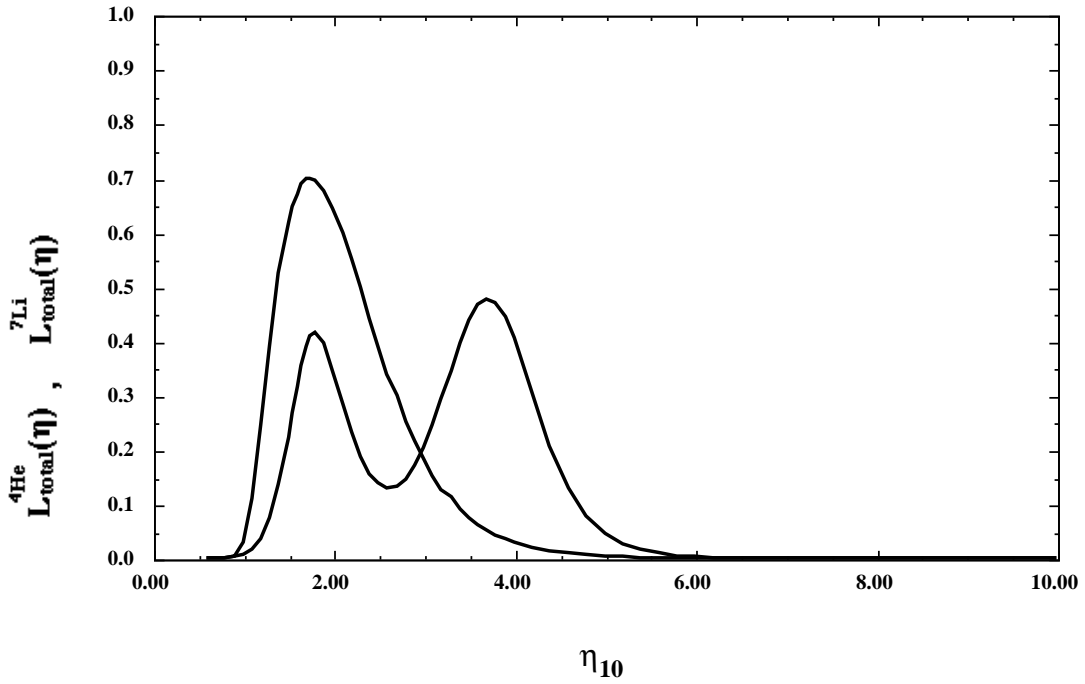
Figure 2: Likelihood distribution for each of ^{4}He and ^{7}Li, shown as a function of η. The one-peak structure of the ^{4}He curve corresponds to its monotonic increase with η, while the two-peaks for ^{7}Li arise from its passing through a minimum.

two error functions

$$L_{\mathrm{O}}(Y, Y_{\mathrm{O}}) = \mathrm{erf}\left(\frac{Y - Y_{\mathrm{O}} + \sigma_{\mathrm{sys}}}{\sqrt{2}\sigma_2}\right) - \mathrm{erf}\left(\frac{Y - Y_{\mathrm{O}} - \sigma_{\mathrm{sys}}}{\sqrt{2}\sigma_2}\right) \tag{25}$$

where in this case, Y_{O} is the observed (or observationally determined) value for the ^{4}He mass fraction. (Had I used a Gaussian to describe the systematic uncertainty, the convolution of two Gaussians leads to a Gaussian, and the likelihood function (25) would have taken a form similar to that in (24).

A total likelihood function for each value of η_{10} is derived by convolving the theoretical and observational distributions, which for ^{4}He is given by

$$L^{^4\mathrm{He}}_{\mathrm{total}}(\eta) = \int dY\, L_{\mathrm{BBN}}\left(Y, Y_{\mathrm{BBN}}\left(\eta\right)\right) L_{\mathrm{O}}(Y, Y_{\mathrm{O}}) \tag{26}$$

An analogous calculation is performed [18] for ^{7}Li. The resulting likelihood functions from the observed abundances given in Eqs. (14) and (15) is shown in Figure 2. As one can see there is very good agreement between ^{4}He and ^{7}Li in the vicinity of $\eta_{10} \simeq 1.8$.

The combined likelihood, for fitting both elements simultaneously, is given by the product of the two functions in Figure 2 and is shown in Figure 3. From Figure 2 it is clear that ^{4}He overlaps the lower (in η) ^{7}Li peak, and so one expects that there will be concordance in an allowed range of η given by the overlap region. This is what one finds in Figure 3, which does show concordance and gives a preferred value for η, $\eta_{10} = 1.8^{+1}_{-.2}$ corresponding to $\Omega_B h^2 = .006^{+.004}_{-.001}$.

Thus, we can conclude that the abundances of ^{4}He and ^{7}Li are consistent, and select an η_{10} range which overlaps with (at the 95% CL) the longstanding favorite range around $\eta_{10} = 3$. Furthermore, by finding concordance using only ^{4}He and ^{7}Li, we deduce that if there

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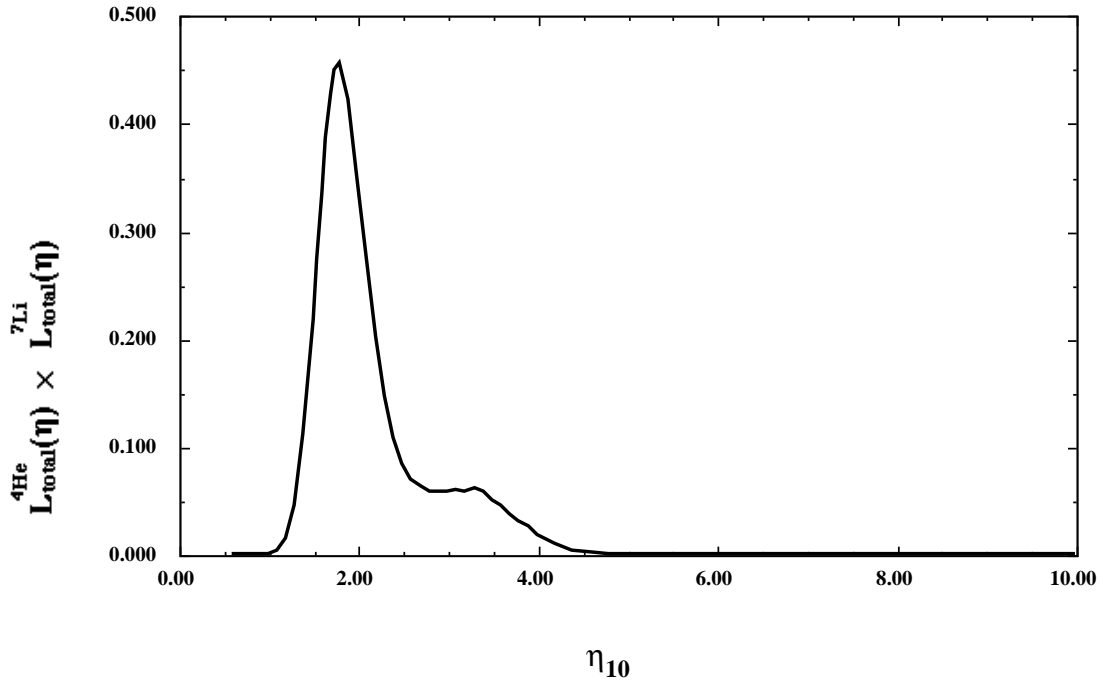Figure 3: Combined likelihood for simultaneously fitting ^{4}He and ^{7}Li, as a function of η.

is problem with BBN, it must arise from D and ^{3}He and is thus tied to chemical evolution or the stellar evolution of ^{3}He. The most model-independent conclusion is that standard BBN with $N_\nu = 3$ is not in jeopardy, but there may be problems with our detailed understanding of D and particularly ^{3}He chemical evolution. It is interesting to note that the central (and strongly) peaked value of η_{10} determined from the combined ^{4}He and ^{7}Li likelihoods is at $\eta_{10} = 1.8$. The corresponding value of D/H is 1.8×10^{-4}, very close to the high value of D/H in quasar absorbers [51, 52, 54]. Since D and ^{3}He are monotonic functions of η, a prediction for η, based on ^{4}He and ^{7}Li, can be turned into a prediction for D and ^{3}He. The corresponding 95% CL ranges are D/H $= (5.5 - 27) \times 10^{-5}$ and and ^{3}He/H $= (1.4 - 2.7) \times 10^{-5}$.

If we did have full confidence in the measured value of D/H in quasar absorption systems, then we could perform the same statistical analysis using ^{4}He, ^{7}Li, and D. To include D/H, one would proceed in much the same way as with the other two light elements. We compute likelihood functions for the BBN predictions as in Eq. (24) and the likelihood function for the observations using D/H $= (1.9 \pm 0.4) \times 10^{-4}$. We are using only the high value of D/H here. These are then convolved as in Eq. (26). In figure 4, the resulting normalized distribution, $L_{total}^D(\eta)$ is super-imposed on distributions appearing in figure 2. It is indeed startling how the three peaks, for D, ^{4}He and ^{7}Li are literally on top of each other. In figure 5, the combined distribution is shown. We now have a very clean distribution and prediction for η, $\eta_{10} = 1.75^{+.3}_{-.1}$ corresponding to $\Omega_B h^2 = .006^{+.001}_{-.0004}$, with the peak of the distribution at $\eta_{10} = 1.75$. The absence of any overlap with the high-η peak of the ^{7}Li distribution has considerably lowered the upper limit to η. Overall, the concordance limits in this case are dominated by the deuterium likelihood function.

To summarize on the subject of big bang nucleosynthesis, I would assert that one can conclude that the present data on the abundances of the light element isotopes are consistent with the standard model of big bang nucleosynthesis. Using the the isotopes with the best data, ^{4}He and ^{7}Li, it is possible to constrain the theory and obtain a best value for the baryon-
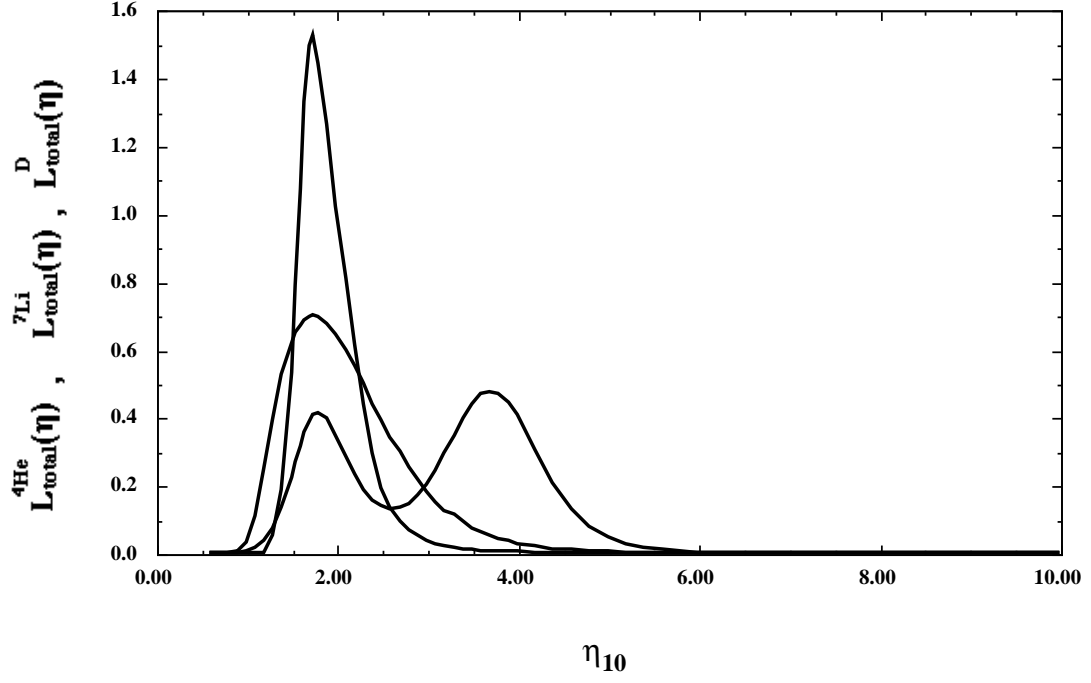
Figure 4: As in Figure 2, with the addition of the likelihood distribution for D/H.
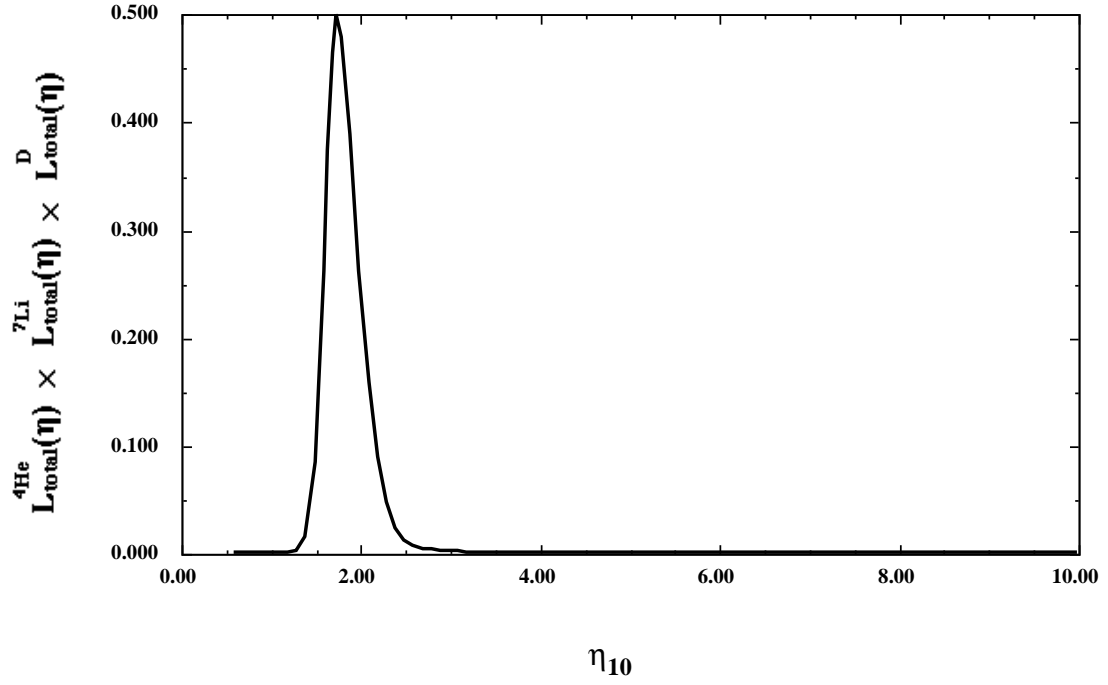

Figure 5: Combined likelihood for simultaneously fitting ^{4}He and ^{7}Li, and D as a function of η.

to-photon ratio of $\eta_{10} = 1.8$, a corresponding value $\Omega_B = .0065$ and

$$
\begin{aligned}
1.4 \quad &< \eta_{10} \quad < 3.8 \qquad 95\%\text{CL} \\
.005 \quad &< \Omega_B h^2 \quad < .014 \qquad 95\%\text{CL}
\end{aligned} \tag{27}
$$

For $0.4 < h < 1$, we have a range $.005 < \Omega_B < .09$. This is a rather low value for the baryon density and would suggest that much of the galactic dark matter is non-baryonic [71]. These predictions are in addition consistent with recent observations of D/H in quasar absorption systems which show a high value. Difficulty remains however, in matching the solar ^{3}He abundance, suggesting a problem with our current understanding of galactic chemical evolution or the stellar evolution of low mass stars as they pertain to ^{3}He.

If we now take as given that non-baryonic dark matter is required, we are faced with the problem of its identity. Light neutrinos ($m \leq 30eV$) are a long-time standard when it comes to non-baryonic dark matter [72]. Light neutrinos produce structure on large scales and the natural (minimal) scale for structure clustering is given in Eq. (11). Hence neutrinos offer the natural possibility for large scale structures [73, 74] including filaments and voids. It seemed, however, that neutrinos were ruled out because they tend to produce too much large scale structure [75]. Because the smallest non-linear structures have mass scale M_J and the typical galactic mass scale is $\simeq 10^{12} M_\odot$, galaxies must fragment out of the larger pancake-like objects. The problem is that in such a scenario, galaxies form late [74, 76] ($z \leq 1$) whereas quasars and galaxies are seen out to redshifts $z \gtrsim 4$. Recently, neutrinos are seeing somewhat of a revival in popularity in mixed dark matter models.

In the standard model, the absence of a right-handed neutrino state precludes the existence of a neutrino mass. By adding a right-handed state ν_R, it is possible to generate a Dirac mass for the neutrino, $m_\nu = h_\nu v/\sqrt{2}$, as for the charged lepton masses, where h_ν is the neutrino Yukawa coupling constant, and v is the Higgs expectation value. It is also possible to generate a Majorana mass for the neutrino when in addition to the Dirac mass term, $m_\nu \bar{\nu}_R \nu_L$, a term $M \nu_R \nu_R$ is included. In what is known as the see-saw mechanism, the two mass eigenstates are given by $m_{\nu_1} \sim m_\nu^2/M$ which is very light, and $m_{\nu_2} \sim M$ which is heavy. The state ν_1 is our hot dark matter candidate as ν_2 is in general not stable.

The cosmological constraint on the mass of a light neutrino is derived from the overall mass density of the Universe. In general, the mass density of a light particle χ can be expressed as

$$
\rho_\chi = m_\chi Y_\chi n_\gamma \leq \rho_c = 1.06 \times 10^{-5} h_o^2 \text{GeV/cm}^3 \tag{28}
$$

where $Y_\chi = n_\chi/n_\gamma$ is the density of χ's relative to the density of photons, for $\Omega h_o^2 < 1$. For neutrinos $Y_\nu = 3/11$, and one finds [77]

$$
\sum_\nu (\frac{g_\nu}{2}) m_\nu < 93\text{eV}(\Omega h_o^2) \tag{29}
$$

where the sum runs over neutrino flavors. All particles with abundances Y similar to neutrinos will have a mass limit given in Eq. (29).

It was possible that neutrinos (though not any of the known flavors) could have had large masses, $m_\nu > 1$ MeV. In that case their abundance Y is controlled by $\nu, \bar{\nu}$ annihilations [78], for example, $\nu\bar{\nu} \to f\bar{f}$ via Z exchange. When the annihilations freeze-out (the annihilation rate becomes slower than the expansion rate of the Universe), Y becomes fixed. Roughly, $Y \sim (m\sigma_A)^{-1}$ and $\rho \sim \sigma_A^{-1}$ where σ_A is the annihilation cross-section. For neutrinos, we

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expect $\sigma_A \sim m_\nu{}^2/m_Z{}^4$ so that $\rho_\nu \sim 1/m_\nu{}^2$ and we can derive a lower bound [79, 80] on the neutrino mass, $m_\nu \gtrsim 3 - 7$ GeV, depending on whether it is a Dirac or Majorana neutrino. Indeed, any particle with roughly a weak scale cross-sections will tend to give an interesting value of $\Omega h^2 \sim 1$ [81].

Due primarily to the limits from LEP [82], the heavy massive neutrino has become simply an example and is no longer a dark matter candidate. LEP excludes neutrinos (with standard weak interactions) with masses $m_\nu \lesssim 40$ GeV. Lab constraints for Dirac neutrinos are available [83], excluding neutrinos with masses between 10 GeV and 4.7 TeV. This is significant, since it precludes the possibility of neutrino dark matter based on an asymmetry between ν and $\bar{\nu}$ [84]. Majorana neutrinos are excluded as *dark matter* since $\Omega_\nu h_o{}^2 < 0.001$ for $m_\nu > 40$ GeV and are thus cosmologically uninteresting.

Supersymmetric theories introduce several possible candidates. If R-parity, which distinguishes between "normal" matter and the supersymmetric partners and can be defined in terms of baryon, lepton and spin as $R = (-1)^{3B+L+2S}$, is unbroken, there is at least one supersymmetric particle which must be stable. I will assume R-parity conservation, which is common in the MSSM. R-parity is generally assumed in order to justify the absence of superpotential terms can be responsible for rampid proton decay. The stable particle (usually called the LSP) is most probably some linear combination of the only $R = -1$ neutral fermions, the neutralinos [85]: the wino $\tilde{W}^3$, the partner of the 3rd component of the $SU(2)_L$ gauge boson; the bino, $\tilde{B}$, the partner of the $U(1)_Y$ gauge boson; and the two neutral Higgsinos, $\tilde{H}_1$ and $\tilde{H}_2$. Gluinos are expected to be heavier -$m_{\tilde{g}} = (\frac{\alpha_3}{\alpha}) \sin^2 \theta_W M_2$, where M_2 is the supersymmetry breaking SU(2) gaugino mass- and they do not mix with the other states. The sneutrino [86] is also a possibility but has been excluded as a dark matter candidate by direct [83] searches, indirect [87] and accelerator[82] searches. For more on the motivations for supersymmetry and the supersymmetric parameter space, see the contribution of Jungman [81].

The identity of the LSP is effectively determined by three parameter in the MSSM, the gaugino mass, M_2, the Higgs mixing mass μ, and the ratio of the Higgs vacuum expectation values, $\tan \beta$. In Figure 6 [89], regions in the M_2, μ plane with $\tan \beta = 2$ are shown in which the LSP is one of several nearly pure states, the photino, $\tilde{\gamma}$, the U(1) gaugino, $\tilde{B}$, a symmetric combination of the Higgsinos, $\tilde{H}_{(12)} = \frac{1}{\sqrt{2}}(\tilde{H}_1 + \tilde{H}_2)$, or the Higgsino $\tilde{S} = \tilde{H}_1 \cos \beta + \tilde{H}_2 \sin \beta$. The dashed lines show the LSP mass contours. The cross hatched regions correspond to parameters giving a chargino ($\tilde{W}^\pm, \tilde{H}^\pm$) state with mass $m_{\tilde{\chi}} \leq 45 GeV$ and as such are excluded by LEP[90]. This constraint has been extended by LEP1.5, [91, 92] and is shown by the light shaded region and corresponds to regions where the chargino mass is $\lesssim 67$ GeV. The dark shaded region corresponds to a limit on M_2 from the limit[93] on the gluino mass $m_{\tilde{g}} \leq 70$ GeV or $M_2 \leq 22$ GeV. Notice that the parameter space is dominated by the $\tilde{B}$ or $\tilde{H}_{12}$ pure states and that the photino (most often quoted as the LSP) only occupies a small fraction of the parameter space, as does the Higgsino combination $\tilde{S}^0$.

As described in [81], the relic abundance of LSP's is determined by solving the Boltzmann equation for the LSP number density in an expanding Universe. The technique[80] used is similar to that for computing the relic abundance of massive neutrinos[78]. For binos, as was the case for photinos [88], it is possible to adjust the sfermion masses $m_{\tilde{f}}$ to obtain closure density. Adjusting the sfermion mixing parameters [95] or CP violating phases [96] allows even greater freedom. In Figure 7 [97], the relic abundance (Ωh^2) is shown in the $M_2 - \mu$ plane with $\tan \beta = 2$, the Higgs pseudoscalar mass $m_0 = 50$ GeV, $m_t = 100$ GeV, and $m_{\tilde{f}} = 200$ GeV. Clearly the MSSM offers sufficient room to solve the dark matter problem. Similar results have
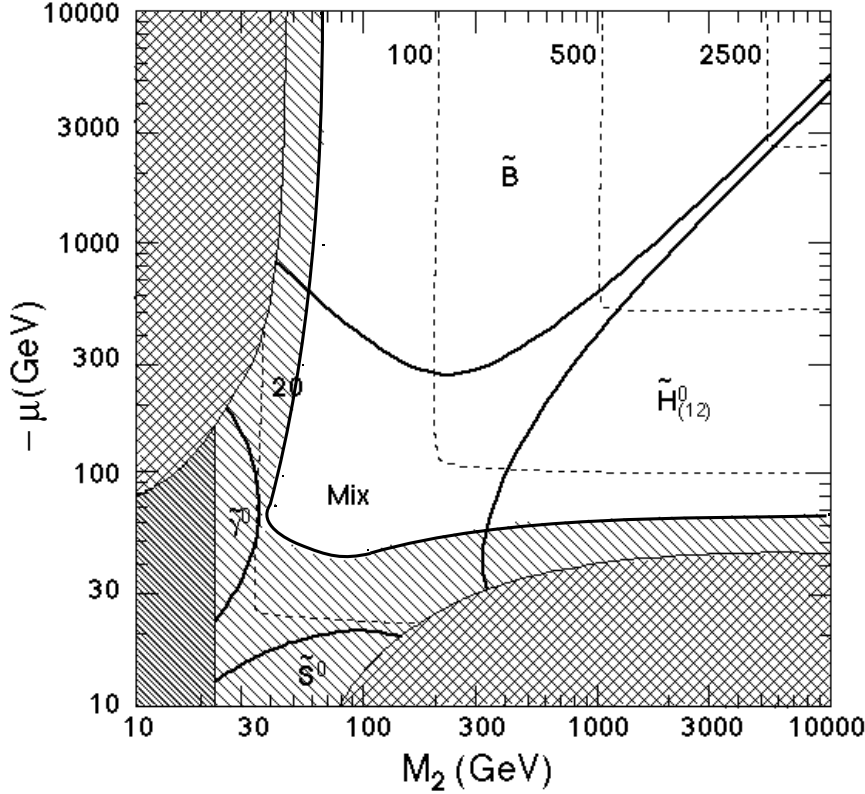
Figure 6: The M_2-μ plane in the MSSM for $\tan\beta = 2$.

been found by other groups [98, 99, 100]. In Figure 7, in the higgsino sector $\tilde{H}_{12}$ marked off by the dashed line, co-annihilations [101, 99] between $\tilde{H}_{(12)}$ and the next lightest neutralino (also a Higgsino) were not included. These tend to lower significantly the relic abundance in much of this sector.

Though I have concentrated on the LSP in the MSSM as a cold dark matter candidate, there are many other possibilities when one goes beyond the standard model. Axions were discussed at length by Jungman [81] and Lu [102]. A host of other possibilities were discussed by Khlopov [103].

The final subject that I will cover in this introduction/summary is the question of detection. Dark matter detection can be separated into two basic methods, direct [104] and indirect [105]. Direct detection relies on the ability to detect the elastic scattering of a dark matter candidate off a nucleus in a detector. The experimental signatures for direct detection were covered by Cabrera [106] and several individual experiments were described [107].

The detection rate, will depend on the density of dark matter in the solar neighborhood, $\rho \sim 0.3$ GeV/cm^3, the velocity, $v \sim 300$ km/s, and the elastic cross section, σ. Spin independent interactions are the most promising for detection. Dirac neutrinos have spin-independent interactions, but as noted above, these have already been excluded as dark matter by direct detection experiments [83]. In the MSSM, it is possible for the LSP to also have spin independent interactions which are mediated by Higgs exchange. These scatterings are only important when the LSP is a mixed (gaugino/Higgsino) state as in the central regions of Figures 6 and 7. Generally, these regions have low values of Ωh^2 (since the annihilation cross sections are also enhanced) and the parameter space in which the elastic cross section and relic density are large
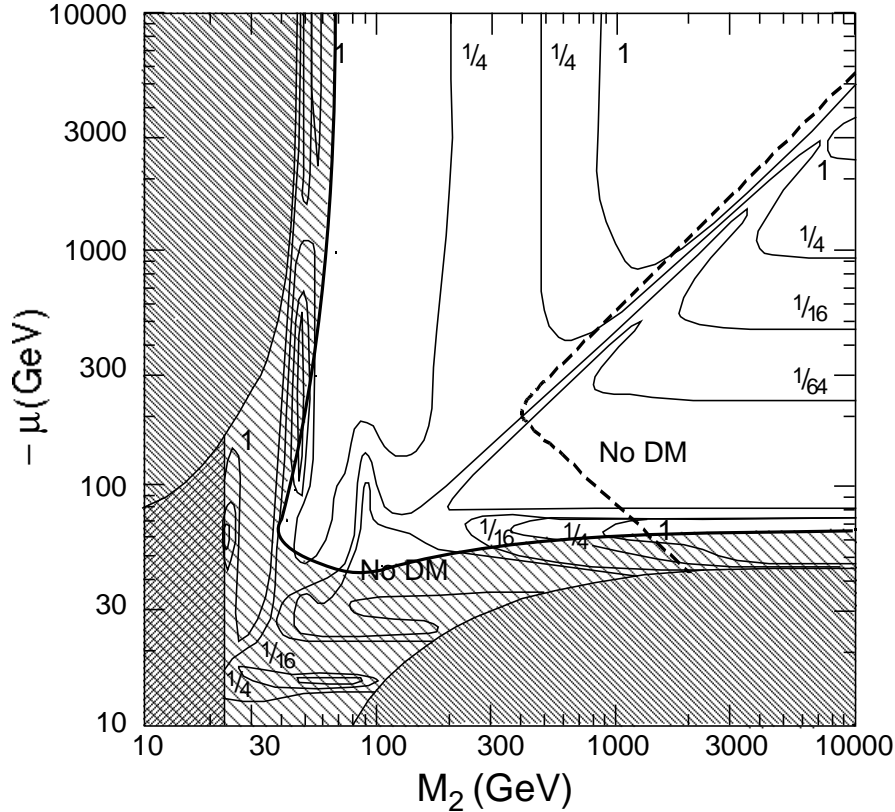
Figure 7: Relic density contours (Ωh^2) in the M_2 - μ plane.

is rather limited. Furthermore, a significant detection rate in this case relies on a low mass for the Higgs scalar [108, 109].

More typical of the SUSY parameter space is a LSP with spin dependent interactions. Elastic scatterings are primarily spin dependent whenever the LSP is mostly either gaugino or Higgsino. For Higgsino dark matter, Higgsinos with scatterings mediated by Z^0 avoid the $\tilde{H}_{(12)}$ regions of Figures 6 and 7, and as such are now largely excluded (the S^0 region does grow at low $\tan\beta$ [85, 89]. Higgsino scatterings mediated by sfermion exchange depend on couplings proportional to the light quark masses and will have cross sections which are suppressed by $(m_p/m_W)^4$, where m_p is the proton mass. These rates are generally very low [109]. Binos, on the other hand, will have elastic cross sections which go as $m^2/m_{\tilde{f}}^4$, where m is the reduced mass of the bino and nucleus. These rates are typically higher (reaching up to almost 0.1 events per kg-day [109, 110, 111].

Indirect methods also offer the possibility for the detection of dark matter. Three methods for indirect detection were discussed [105]. 1) γ-rays from dark matter annihilations in the galactic halo are a possible signature [112]. In the case of the MSSM, unless the mass of the LSP is larger than m_W, the rates are probably too small to be detectable over background [105]. 2) Dark matter will be trapped gradually in the sun, and annihilations within the sun will produce high energy neutrinos which may be detected [113]; similarly, annihilations within the earth may provide a detectable neutrino signal [114]. Edsjo [115] discussed possibilities for determining the mass the dark matter candidate from the angular distribution of neutrinos. This method hold considerable promise, as there will be a number of very large neutrino detectors coming on line in the future. Finally, 3) there is the possibility that halo annihilations into positrons

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and antiprotons in sufficient numbers to distinguish them from cosmic-ray backgrounds [112, 116, 117].

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